

# Constraints on both bilinear and trilinear R-parity violating couplings from neutrino laboratories and astrophysics data

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## Abstract

We consider neutrino masses generated at tree level and at one loop, through fermion–sfermion loop diagrams, in the MSSM with R-parity violation. Using the  $(3 \times 3)$  mass and mixing matrices for three generations of neutrinos and the present experimental results on neutrinos from laboratories and astrophysics simultaneously, we put bounds on both trilinear  $(\lambda_{ijk}, \lambda'_{ijk})$  and bilinear  $(\mu_e, \mu_\mu, \mu_\tau)$  R-parity-violating couplings.

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## I. INTRODUCTION

The analysis of the flavour structure of the neutrino mass matrix has been a subject of intensive study recently [1]. A non-trivial neutrino mass matrix can solve, through the oscillation solutions, the atmospheric and solar neutrino anomalies, which have been observed by different experiments: to the very recent results of the Super-Kamiokande collaboration [2], have to be added those of other atmospheric neutrino experiments (IMB [3], Soudan [4], Kamiokande [5]) and solar neutrino experiments (Homestake [6], Gallex [7], SAGE [8], Kamiokande [9], Super-Kamiokande [10], MACRO [11] and LSND [12]).

The oscillation explanation of the solar neutrino problem, the atmospheric neutrino anomaly and the LSND results suggest three very different values of neutrino mass squared differences, namely  $\Delta m_{\text{sun}}^2 \ll \Delta m_{\text{atm}}^2 \ll \Delta m_{\text{LSND}}^2$ , with  $\Delta m_{\text{sun}}^2 \lesssim 10^{-4} \text{ eV}^2$ ,  $\Delta m_{\text{atm}}^2 \in [10^{-3}, 10^{-2}] \text{ eV}^2$  and  $\Delta m_{\text{LSND}}^2 \in [0.3, 1.0] \text{ eV}^2$ . The evidence for  $\bar{\nu}_\mu \rightarrow \bar{\nu}_e$  observed by LSND has not been confirmed or excluded by the KARMEN [13] experiment. MiniBooNE at FNAL [14] or MINOS long-baseline experiments [15] could provide the answer. Excluding LSND results, the oscillation explanation of the solar neutrino problem and the atmospheric neutrino anomaly requires two mixing angles and two values of neutrino mass squared differences with a strong hierarchy, namely  $\Delta m_{\text{sun}}^2 \ll \Delta m_{\text{atm}}^2$ .

In a previous analysis [16] with a general  $(3 \times 3)$  symmetric mass matrix assuming oscillations as an explanation of the neutrino deficits, we had used as inputs neutrino data constraints in order to obtain the allowed range of variation of the mass matrix elements. These inputs are  $\Delta m_{\text{atm}}^2, \sin^2 2\theta_{\text{atm}}$  from SuperK measurements,  $\Delta m_{\text{sun}}^2, \sin^2 2\theta_{\text{sun}}$  from solar experiments and finally the CHOOZ constraint [17]  $\sin^2 \theta_{\text{CHOOZ}}$ . We also discussed the spectrum after diagonalization and considered two interesting quantities: the sum of the neutrino masses, which is a relevant quantity for obtaining a neutrino component of hot dark matter (HDM), and the effective mass constraint, which is relevant for neutrinoless double beta decay (this is applicable when the neutrinos are Majorana and massive particles).

To generate a Majorana mass for the neutrinos in a given model lepton number must be violated. We consider the Minimal Supersymmetric Standard Model (MSSM) with R-parity [18] violation ( $R_P = (-1)^{L+3B+2S}$ , where  $L$ ,  $B$ ,  $S$  are the lepton and baryon number and the spin of the particle, respectively), left-handed neutrinos obtain a Majorana mass, at tree level, through mixing with the neutralinos, and through loop diagrams that violate lepton number (in two units). The simultaneous presence of baryon and lepton number violating couplings is not acceptable, because of the long lifetime of the proton. As we are interested in a model that provides lepton-number violation, we simply choose all baryon number violating couplings to be exactly zero.

In ref. [16] we presented as an application of our results the case of the Minimal Supersymmetric Standard Model (MSSM), with R-parity violation at the one-loop order allowing

the presence of both bilinear and trilinear  $R_P$ -violating couplings. We included the effect of the bilinear couplings at tree level and of trilinear couplings at one loop. The numerical bounds were given assuming all trilinear  $\lambda_{ijk} \equiv \lambda$ , as well as  $\lambda'_{ijk} \equiv \lambda'$  to be equal. Here we will relax this assumption and apply our general results to other limiting (albeit general) cases in this model. The aim is to provide strong constraints on the couplings with an explicit reference to the generation indices from neutrino data. We will allow the simultaneous presence of various lepton number violating couplings.

The paper is organized as follows. In section II we describe the inputs used to constrain the  $R$ -parity violating parameters, while in section III we introduce our notation for the MSSM with  $R_P$ -violation. We also give in section III the tree-level neutrino  $(3 \times 3)$  mass matrix and the loop corrections to each matrix element. We parametrize the mass and mixing matrices. In section IV, we give results for five separate subcases that correspond to some limiting cases of the MSSM without  $R$ -parity, and we derive bounds on the bilinear and trilinear  $R_P$ -violating couplings for each case.

## II. GENERAL INPUTS

Global fits of neutrino data have shown that neutrino oscillations among three flavours are sufficient to accommodate the solar and atmospheric data [19]. Three different types of solar experiments are sensitive to different solar neutrino energy ranges; consequently, three different ranges of solutions for the solar data exist, which correspond to the vacuum oscillation solution, MSW with a large mixing angle (MSW-LMA) and MSW with a small mixing angle (MSW-SMA). The required neutrino mass squared differences and mixing angles are shown in table I. We use these inputs together with the results from the CHOOZ experiment [17] to constrain the elements of our neutrino mass matrix in the flavour basis. An additional sixth input is needed to solve the general  $(3 \times 3)$  real and symmetric mass matrix. As our sixth input we could take the direct upper bound on the effective mass from neutrinoless beta decay  $(\beta\beta)_{0\nu}$ , or the upper bound on the effective neutrino mass coming from the  $^3\text{H}$   $\beta$ -decay spectrum, or the astrophysical bound on the magnetic moment of the neutrino bounds, or the global cosmological upper bound on the sum of the neutrino masses [20].

Let us briefly discuss each of these possible inputs separately. The direct upper bound from the measurements of the high energy part of the  $^3\text{H}$   $\beta$ -decay spectrum given as upper limits on the electron neutrino mass obtained in the TROITSK [21] and MAINZ [22] experiments are  $m_\nu < 2.5$  eV and  $m_\nu < 2.8$  eV, respectively. However, these experiments suffer from some ambiguities referred to as “the negative mass squared problem”, which is still not completely understood and we will therefore not use these bounds.

The upper bound on the sum of the neutrino masses coming from astrophysical and cosmological considerations, such as the one from hot dark matter (HDM), suggests that

$$\sum_i m_{\nu_i} = m_1 + m_2 + m_3 < \text{few eV},$$

where  $m_i, i = 1, 2, 3$ , are the masses of the mass eigenstates  $\nu_i$ , which constitute active flavour neutrinos. A recent analysis, based on the observation of distant objects favouring a non-zero cosmological constant instead of HDM [23], shows that the bound of a few eV for the sum of the neutrino masses is no longer required, since a HDM component is not a necessary ingredient in this case.

In the R-parity violating model that we are considering we can have transition magnetic moments that change lepton number. These magnetic moments are proportional to the matrix elements; we prefer to use as an input a direct bound on a matrix element as such arising from neutrinoless double beta decay.

The additional input we will take is the best limit on the effective mass  $m_{\text{eff}}$  appearing in  $(\beta\beta)_{0\nu}$ , as it directly constrains the matrix element  $m_{11}$  of the neutrino mass matrix. It is defined by

$$|m_{\text{eff}}| = \left| \sum_i m_{\nu_i} U_{ei}^2 \right| \leq \sum_k m_{\nu_k} |U_{ek}^2|, \quad (1)$$

which has been derived in the Heidelberg–Moscow  $^{76}\text{Ge}$  experiment [24] (see also references [25,26]):

$$|m_{\text{eff}}| < (0.2 - 0.6) \text{ eV} \quad (\text{at } 90\% \text{CL}). \quad (2)$$

The CP phases that might appear in the mixing matrix elements  $U_{ei}$  are not relevant to our analysis, as we only use the second r.h.s. term in eq. (1).

For three generations, the flavour states  $\nu_l$  are expressed in terms of the mass eigenstates  $\nu_i$  using the  $(3 \times 3)$  mixing matrix  $U$

$$\nu_\ell = \sum_{i=1}^3 U_{\ell i}^* \nu_i. \quad (3)$$

Using the Chau and Keung parametrization of a  $3 \times 3$  rotation matrix [27],  $U$  is given by

$$U = \begin{pmatrix} c_{12} c_{13} & s_{12} c_{13} & s_{13} e^{-i\delta} \\ -s_{12} c_{23} - c_{12} s_{23} s_{13} e^{i\delta} & c_{12} c_{23} - s_{12} s_{23} s_{13} e^{i\delta} & s_{23} c_{13} \\ s_{12} s_{23} - c_{12} c_{23} s_{13} e^{i\delta} & -c_{12} s_{23} - s_{12} c_{23} s_{13} e^{i\delta} & c_{23} c_{13} \end{pmatrix} \text{diag} \{ e^{i\alpha_1}, e^{i\alpha_2}, 1 \}, \quad (4)$$

where  $c_{ij} = \cos(\theta_{ij})$  and  $s_{ij} = \sin(\theta_{ij})$ ,  $\delta$  is the Dirac CP phase and  $\alpha_{1,2}$  are the Majorana ones (we have two additional CP phases in the case of Majorana particles).

The survival probability  $P_{\nu_e \rightarrow \nu_e}$ , relevant to the case of solar fluxes, depends only on the first row of the mixing matrix in eq. (4), i.e. on  $|U_{ei}|^2$ , with  $i = 1, 2, 3$ . In the atmospheric case, the oscillation probability depends on the last column of (4), i.e. on  $|U_{\ell 3}|^2$ , with  $\ell = e, \mu, \tau$ . The other elements of the matrix are not constrained by any direct experimental observation. With this parametrization, we directly obtain the mixing angles  $\sin^2 2\theta_{12}$  and  $\sin^2 2\theta_{23}$ , and also the relevant CHOOZ parameter  $\sin \theta_{13}$ .

We impose the following hierarchy on the mass eigenvalues and denote them by  $m_i$ ,  $i = 1, 2, 3$ , such that  $m_1 \leq m_2 \leq m_3$ . As the neutrino oscillation scenario cannot fix the absolute mass scale nor distinguish whether the smallest mass splitting is between the two lightest mass eigenstates or the two heaviest ones, we present in table II the possible neutrino spectra and indicate which is the corresponding mass squared difference.

To summarize, we will use the following six inputs from neutrino data:

- the limit from the effective mass appearing in neutrinoless double beta decay:  $m_{\text{eff}}$ ,
- $\Delta m_{\text{atm}}^2$  and  $\sin^2 2\theta_{\text{atm}}$ ,
- $\Delta m_{\text{sun}}^2$  and  $\sin^2 2\theta_{\text{sun}}$ ,
- $\sin^2 2\theta_{\text{CHOOZ}}$ .

Experiment	$\Delta m^2$ (eV <sup>2</sup> )	$\sin^2 2\theta$
Atmospheric	$(2 - 5) \times 10^{-3}$	$0.88 - 1$
Solar		
MSW-LMA	$(3 - 30) \times 10^{-5}$	$0.6 - 1$
MSW-SMA	$(0.4 - 1) \times 10^{-5}$	$10^{-3} - 10^{-2}$
Vacuum	$(0.5 - 8) \times 10^{-10}$	$0.5 - 1$
CHOOZ	$> 3 \times 10^{-3}$	$< 0.22$

TABLE I. MSW-LMA, MSW-SMA and Vacuum stand for MSW large mixing angle, small mixing and vacuum oscillation solutions, respectively.

Spectrum	Solar	Atmospheric
Hierarchy	$\Delta m_{12}^2$	$\Delta m_{13}^2$
Degenerate	$\Delta m_{23}^2$ or $\Delta m_{12}^2$	$\Delta m_{13}^2$
Pseudo-Dirac	$\Delta m_{23}^2$ or $\Delta m_{12}^2$	$\Delta m_{13}^2$

TABLE II. Different possible regimes and corresponding mass squared difference.

### III. MSSM WITH R<sub>P</sub>- VIOLATION

The most general renormalizable superpotential for the supersymmetric Standard Model with lepton-number violation is

$$W = \epsilon_{ab}[\mu_\alpha \hat{L}_\alpha^a \hat{H}_u^b + \lambda_{\alpha\beta k} L_\alpha^a \hat{L}_\beta^b \hat{E}_k^C + h_{ik}^u \hat{Q}_i^a \hat{H}_u^b \hat{U}_k^C + \lambda'_{\alpha ik} \hat{L}_\alpha^b \hat{Q}_i^a \hat{D}_k^C], \quad (5)$$

where the  $(i, j, k)$  are flavour indices,  $(a, b)$  are  $SU(2)$  indices, and the  $(\alpha, \beta)$  are flavour indices running from 0 to 3. The  $\hat{L}_\alpha$  are the doublet superfields with hypercharge  $Y = -1$ . Note that the  $\lambda$  couplings are antisymmetric in the first two indices. The usual R-parity-preserving Lagrangian is obtained when only  $\mu_o, \lambda_{0ik} = h_{ik}^d, \lambda'_{i0k} = h_{ik}^d$  are non-zero, and we can identify  $\hat{L}_o \equiv \hat{H}_d$ . In the model of eq. (5) we have 9 additional  $\lambda$  couplings and 27 new  $\lambda'$  couplings with respect to the R-parity-conserving case. Note that thanks to the additional degrees of freedom, we can rotate in the flavour space of the “down-type” scalar fields to set the vacuum expectation values of the sneutrinos to be zero<sup>1</sup>. Henceforth, we will work consistently in this basis and the bounds we will derive are valid for this basis.

This model has been extensively analysed in the literature [28]–[31]. References [30,31] showed in a basis-invariant way that neutrino masses are always generated in these models, even when universality of the soft SUSY-breaking terms is assumed at some high scale. Previous studies have also tried to constrain the different R<sub>P</sub>-violating couplings that appear in the MSSM Lagrangian, considering only the effect of bilinear terms [32] or only of trilinear couplings [33], or of both [34,16], from solar and atmospheric neutrino data. Both tree-level and one-loop effects have been considered. A recent study [35] has constrained these couplings using rare decays. The results are in agreement with the constraints given in ref. [16].

It is well known that the tree-level expression of the neutrino mass matrix elements obtained through the mixing with neutralinos is

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<sup>1</sup>This can be done order by order in the loop expansion when one appropriately defines the mass matrices of the Higgs sector.

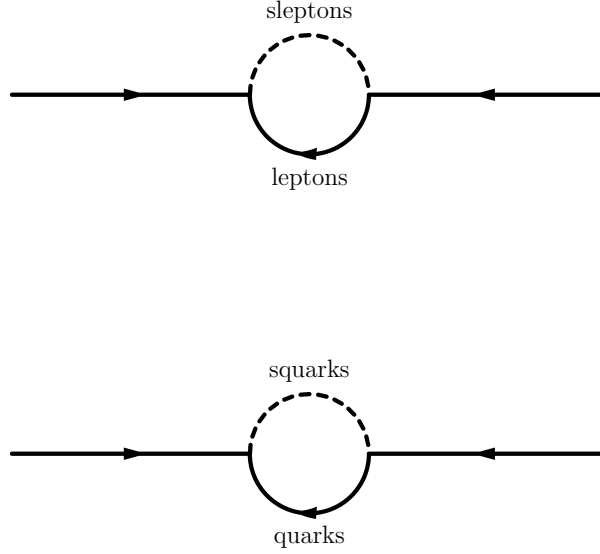


FIG. 1. *One-loop Feynman diagrams contributing to the neutrino masses.*

$$M_{\nu_{ij}}^{\text{tree}} = g_2^2 \frac{(M_1 + \tan^2 \theta_W M_2)}{4 \det M} \mu_i \mu_j v_1^2 \sim g_2^2 \frac{v^2 \cos^2 \beta}{M_{\text{susy}}^3} \mu_i \mu_j \equiv C \mu_i \mu_j, \quad (6)$$

where  $i, j$  are flavour indices,  $M_1, M_2$  are the gaugino masses and  $\det M$  is the determinant of the R-parity conserving neutralino mass matrix. For the second term of the right-hand side we have taken all R-parity conserving masses to be of order  $\sim M_{\text{susy}}$  and as we are working in the basis where the slepton vacuum expectation value is zero,  $v_1 = v_d = v \cos \beta$ . We can rewrite the tree-level mass matrix as

$$\mathcal{M}_\nu^{\text{tree}} = C \begin{pmatrix} \mu_e^2 & \mu_\mu \mu_e & \mu_e \mu_\tau \\ \mu_\mu \mu_e & \mu_\mu^2 & \mu_\mu \mu_\tau \\ \mu_e \mu_\tau & \mu_\mu \mu_\tau & \mu_\tau^2 \end{pmatrix}. \quad (7)$$

The above mass matrix has only one non-zero eigenvalue. In order for the other neutrinos to obtain a mass, we must consider loop corrections. There are many types of loop diagrams that can contribute to the neutrino mass in this model [29,36,37]. Here we will focus on the contribution at one loop from fermion–sfermion diagrams involving trilinear couplings, see fig. 1. See, for example, ref. [16] for more details regarding the computation of the neutrino mass matrix.

The one-loop mass contribution from slepton–lepton and squark–quark loops to each element of the mass matrix is given by

$$(m_{qm})_{\text{loop}} = \frac{1}{16\pi^2} X \left( \frac{f(x_\ell)}{M_2^2} \sum_{k,p} \lambda_{qkp} \lambda_{mpk} m_\ell^{(k)} m_\ell^{(p)} + 3 \frac{f(x_q)}{M_{q2}^2} \sum_{k,p} \lambda'_{qkp} \lambda'_{mpk} m_q^{(k)} m_q^{(p)} \right), \quad (8)$$

where

$$f(x) = -\frac{\ln x}{1-x}, \quad x_\ell^{(p)} = \left(\frac{M_1^{(p)}}{M_2^{(p)}}\right)^2, \quad x_q^{(p)} = \left(\frac{M_{q1}^{(p)}}{M_{q2}^{(p)}}\right)^2 \quad \text{and} \quad X = A + \mu \tan \beta, \quad (9)$$

where  $X$  is the trilinear term that appears in the off-diagonal matrix element of the sfermions mass matrix. We consider here that we are in the down-quark mass eigenstate basis, and that the  $\lambda, \lambda'$  couplings have been redefined in terms of the couplings appearing in the superpotential and of the corresponding Cabibbo–Kobayashi–Maskawa matrix elements; we also drop the flavour dependence from  $\frac{X}{M_2^{(p)2}} f(x_\ell^{(p)})$  and  $\frac{X}{M_{q2}^{(p)2}} f(x_q^{(p)})$ , and consider them to be universal in the slepton and in the squark sector, respectively. Thus the matrix elements at one loop are given by the sum of eqs. (7) and (8). The analysis we perform can directly constrain terms of the form  $C\mu_i\mu_j$ ,  $K_1^{ij} \sum_{k,p} \lambda_{ikp} \lambda_{jpk}$ ,  $K_2^{ij} \sum_{k,p} \lambda'_{ikp} \lambda'_{jpk}$ , where

$$K_1^{ij} = \frac{X}{16\pi^2} \frac{f(x_\ell)}{M_2^2} (m_i m_j),$$

$$K_2^{ij} = 3 \frac{X}{16\pi^2} \frac{f(x_q)}{M_2^{q2}} (m_i m_j). \quad (10)$$

Note that the quantities  $K_{1,2}^{ij}$  depend only on R-parity conserving parameters. As an application, we focus in the following on a case where we assume that the  $\lambda'_{ijk}$  are of the same order, which allows us to neglect the extra mixed  $\lambda'_{123,132,223,232}$  due to the hierarchy  $m_s m_b \ll m_b^2$  (see eq. (8)). We also assume separately that all  $\lambda_{ijk}$  are of the same order<sup>2</sup>. Thus, using the mass hierarchy  $m_{e,d} \ll m_{\mu,s} \ll m_{\tau,b}$  in  $(m_{qm})_{\text{loop}}$ , simplifies the expressions for the matrix elements, to obtain

$$\mathcal{M}_\nu^{\text{loop}} = \begin{pmatrix} K_1 \lambda_{133}^2 + K_2 \lambda_{133}'^2 & K_1 \lambda_{133} \lambda_{233} + K_2 \lambda_{133}' \lambda_{233}' & K_2 \lambda_{133}' \lambda_{333}' \\ K_1 \lambda_{133} \lambda_{233} + K_2 \lambda_{133}' \lambda_{233}' & K_1 \lambda_{233}^2 + K_2 \lambda_{233}'^2 & K_2 \lambda_{233}' \lambda_{333}' \\ K_2 \lambda_{133}' \lambda_{333}' & K_2 \lambda_{233}' \lambda_{333}' & K_2 \lambda_{333}'^2 \end{pmatrix}, \quad (11)$$

where the coefficients  $K_{1,2}$  are given by:

$$K_1 = \frac{X}{16\pi^2} \frac{f(x_\ell)}{M_2^2} (m_\tau^2),$$

$$K_2 = 3 \frac{X}{16\pi^2} \frac{f(x_q)}{M_2^{q2}} (m_b^2). \quad (12)$$

Thus, we take our one-loop mass matrix to be

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<sup>2</sup>The same procedure as we apply here could be performed if we considered that only one of the R-parity violating trilinear couplings was dominant.



$$\mathcal{M}_\nu = \mathcal{M}_\nu^{\text{tree}} + \mathcal{M}_\nu^{\text{loop}}. \quad (13)$$

At this order, the mass matrix has in general three non-zero eigenmasses. The **8** R-parity violating parameters that we would like to constrain are:

- $\lambda_{133}, \lambda_{233}$
- $\lambda'_{133}, \lambda'_{233}$  and  $\lambda'_{333}$
- $\mu_e, \mu_\mu$  and  $\mu_\tau$ .

Recall that we will use only **6** inputs from solar (2), atmospheric (2), CHOOZ (1) constraints and the bound on the effective neutrino mass appearing in neutrinoless double beta decay.

In our previous analysis [16], denoted case 0 in table III, we considered a toy model leading to three non-zero mass eigenstates. In this toy model we assumed that all trilinear couplings  $\lambda_{ijk}$  were equal and that all  $\lambda'_{ijk}$  couplings were equal, i.e.  $\lambda_{ijk} = \lambda$  and  $\lambda'_{ijk} = \lambda'$ . This leads to five unknowns  $\mu_e, \mu_\tau$ ,  $\lambda$  and  $\lambda'$ , and the system was solved using the five inputs from the solar, atmospheric and CHOOZ constraints.

In this work, we will relax this assumption to more general cases. We consider five different cases that generate three non-zero eigenmasses. These different subcases of the generic matrix of eq. (13), arise when we apply certain conditions on some of the  $L$ -number violating parameters. We do this so as to reduce the number of unknowns to the number of constraints. The subcases we consider are summarized in table 2.

A first case corresponds to the situation where all the  $\lambda'$  couplings are switched off, which happens in the limit in which the squarks decouple. Case 2 occurs when all the  $\lambda$  couplings are switched off, i.e. in the limit in which the sleptons decouple. In case 3, the bilinear couplings are such that  $\mu_e \sim \mu_\mu \ll \mu_\tau$  and thus only one bilinear coupling ( $\mu_\tau$ ) is relevant. In cases 4 (5), the trilinear  $\lambda$  ( $\lambda'$ ) are of the same order, i.e.  $\lambda_{133} \simeq \lambda_{233} = \lambda$  ( $\lambda'_{133} \simeq \lambda'_{233} \simeq \lambda'_{333} = \lambda'$ ) and the bilinears are such that  $\mu_e \sim \mu_\mu \neq \mu_\tau$  for case 4 and  $\mu_e \neq \mu_\mu \neq \mu_\tau$  for case 5. We do not claim to have covered all possibilities, but certainly some of the most straightforward ones.

Cases	$\lambda_{ijk}$	$\lambda'_{ijk}$	$\mu_i$
0	$\lambda_{ijk} = \lambda$	$\lambda'_{ijk} = \lambda'$	$\mu_e, \mu_\mu, \mu_\tau$
1	$\lambda_{133}, \lambda_{233}$	0	$\mu_e, \mu_\mu, \mu_\tau$
2	0	$\lambda'_{133}, \lambda'_{233}, \lambda'_{333}$	$\mu_e, \mu_\mu, \mu_\tau$
3	$\lambda_{133}, \lambda_{233}$	$\lambda'_{133}, \lambda'_{233}, \lambda'_{333}$	$\mu_\tau$
4	$\lambda_{ijk} = \lambda$	$\lambda'_{133}, \lambda'_{233}, \lambda'_{333}$	$\mu_\tau, \mu_e = \mu_\mu$
5	$\lambda_{133}, \lambda_{233}$	$\lambda'_{ijk} = \lambda'$	$\mu_e, \mu_\mu, \mu_\tau$

TABLE III. Different cases of R<sub>p</sub>-violating couplings contributing to the neutrino mass matrix.

#### IV. RESULTS

In our general scan of parameter space, we allow tree-level contributions to either dominate over the loop corrections, or to be on the same order as these, or to be much smaller than the loop terms. The analysis for subcase 0, presented in [16], gave bounds on the  $R_p$  couplings from combinations of constraints from atmospheric and CHOOZ data, together with one of the possible solar neutrino solutions. For our five new subcases we relax the assumptions of subcase 0 to present stringent constraints on bilinear and trilinear lepton-number violating couplings, with specific generation indices in the basis where the sneutrino vacuum expectation value is zero.

The bounds on the couplings are presented in tables for the various combination of constraints from neutrinoless double beta decay, atmospheric, CHOOZ together with vacuum or MSW-SMA or MSW-LMA solutions. A common assumption used to place bounds in this model is to take all R<sub>p</sub>-conserving mass parameters to be of the same order,  $M_{\text{susy}}$ . For the particular case where  $M_{\text{susy}} = 100 \text{ GeV}$ ,  $f(x) \rightarrow 1$  and  $\tan \beta = 2$ , we have

$$K_1 \sim 1.8 \times 10^{-4} \text{GeV}, \quad K_2 \sim 4.7 \times 10^{-3} \text{GeV}, \quad C \sim 5.3 \times 10^{-3} \text{GeV}^{-1}.$$

We use these values to obtain the numerical results presented for cases 1–5 in tables IV–VIII, respectively. Using eqs. (6) and (12) modified bounds can be obtained for other values of the R-parity conserving parameters.

The analysis for all subcases shows that there are many regions of parameter space that can simultaneously accommodate the MSW solution (large or small mixing angle), SuperK and CHOOZ constraints. Subcases 2–4 also provide solutions for the combined constraint of atmospheric data, CHOOZ and the vacuum oscillation solution, while subcases 1 and 5 do not. The results for the latter two subcases are similar to those obtained for subcase 0 treated in [16]. The number of solutions defining the allowed region in the  $(\Delta m^2, \sin^2 2\theta)$

Couplings	MSW-LMA	MSW-SMA	Vacuum
$ \lambda_{133} $	$6.8 \times 10^{-4}$	$1.3 \times 10^{-4}$	-
$ \lambda_{233} $	$5.6 \times 10^{-4}$	$1.9 \times 10^{-4}$	-
$ \mu_e $ (GeV)	$6.1 \times 10^{-5}$	$2.0 \times 10^{-6}$	-
$ \mu_\mu $ (GeV)	$1.2 \times 10^{-4}$	$1.2 \times 10^{-4}$	-
$ \mu_\tau $ (GeV)	$1.2 \times 10^{-4}$	$1.2 \times 10^{-4}$	-

TABLE IV. Bounds on the couplings for subcase 1 that satisfy MSW-LMA or MSW-SMA, SuperK, CHOOZ and neutrinoless double beta decay simultaneously.

Couplings	MSW-LMA	MSW-SMA	Vacuum
$ \lambda'_{133} $	$1.5 \times 10^{-4}$	$4.1 \times 10^{-5}$	$5.4 \times 10^{-5}$
$ \lambda'_{233} $	$1.5 \times 10^{-4}$	$1.5 \times 10^{-4}$	$9.5 \times 10^{-5}$
$ \lambda'_{333} $	$1.5 \times 10^{-4}$	$1.5 \times 10^{-4}$	$1.5 \times 10^{-5}$
$ \mu_e $ (GeV)	$1.2 \times 10^{-4}$	$4.1 \times 10^{-5}$	$5.1 \times 10^{-5}$
$ \mu_\mu $ (GeV)	$1.2 \times 10^{-4}$	$1.2 \times 10^{-4}$	$8.1 \times 10^{-5}$
$ \mu_\tau $ (GeV)	$1.2 \times 10^{-4}$	$1.2 \times 10^{-4}$	$1.2 \times 10^{-4}$

TABLE V. Bounds on the couplings for subcase 2 that satisfy MSW-LMA or MSW-SMA or vacuum, SuperK, CHOOZ and the neutrinoless beta decay constraints simultaneously.

plane is very small for the vacuum solution, larger in the case of MSW-SMA, and it is still larger in the case of MSW-LMA. This is in agreement with the most recent results from the Super-Kamiokande collaboration, presented at the SUSY2K conference [38], stating that the solar neutrinos favour the MSW-LMA solution. There is less room for MSW-SMA and much less still for vacuum solutions. We can see that our bounds are consistent with bounds derived in the literature (see [39–41] for three-neutrino case and [42] for four-neutrino case). We emphasize that some of our subcases present more stringent bounds, and that the combination of constraints with the vacuum solution requires the smallest values for R-parity violating couplings.

To summarize, we have obtained bounds on the bilinear and trilinear R-parity violating couplings with explicit reference to the leptonic indices. The bounds have been obtained from neutrino data, which constrain the neutrino mass matrix that can be constructed in the MSSM with R-parity violation for three generations of neutrinos. We considered the tree-level contribution and the one-loop contribution from fermion–sfermion diagrams with trilinear couplings to the neutrino mass matrix.

Couplings	MSW-LMA	MSW-SMA	Vacuum
$ \lambda'_{133} $	$1.2 \times 10^{-4}$	$6.8 \times 10^{-5}$	$1.4 \times 10^{-5}$
$ \lambda'_{233} $	$1.5 \times 10^{-4}$	$1.5 \times 10^{-4}$	$1.4 \times 10^{-5}$
$ \lambda'_{333} $	$1.5 \times 10^{-4}$	$1.5 \times 10^{-4}$	$1.5 \times 10^{-4}$
$ \lambda_{133} $	$6.2 \times 10^{-4}$	$3.1 \times 10^{-4}$	$1.1 \times 10^{-7}$
$ \lambda_{233} $	$6.2 \times 10^{-4}$	$3.1 \times 10^{-4}$	$1.1 \times 10^{-7}$
$ \mu_\tau $ (GeV)	$1.3 \times 10^{-4}$	$1.3 \times 10^{-4}$	$1.2 \times 10^{-5}$

TABLE VI. Bounds on the couplings for subcase 3 that satisfy MSW-LMA or MSW-SMA or vacuum, SuperK, CHOOZ and the neutrinoless beta decay constraints simultaneously.

Couplings	MSW-LMA	MSW-SMA	Vacuum
$ \lambda'_{133} $	$1.1 \times 10^{-4}$	$3.4 \times 10^{-5}$	$5.4 \times 10^{-5}$
$ \lambda'_{233} $	$1.4 \times 10^{-4}$	$1.4 \times 10^{-4}$	$5.4 \times 10^{-5}$
$ \lambda'_{333} $	$1.5 \times 10^{-4}$	$1.4 \times 10^{-4}$	$1.5 \times 10^{-4}$
$ \lambda $	$4.4 \times 10^{-4}$	$1.9 \times 10^{-4}$	$1.1 \times 10^{-7}$
$ \mu_e \sim \mu_\mu $ (GeV)	$8.1 \times 10^{-5}$	$2.3 \times 10^{-5}$	$4.6 \times 10^{-5}$
$ \mu_\tau $ (GeV)	$1.2 \times 10^{-4}$	$1.3 \times 10^{-4}$	$1.3 \times 10^{-4}$

TABLE VII. Bounds on the couplings for subcase 4 that satisfy MSW-LMA or MSW-SMA or vacuum, SuperK, CHOOZ and the neutrinoless beta decay constraints simultaneously.

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Couplings	MSW-LMA	MSW-SMA	Vacuum
$ \lambda_{133} $	$6.2 \times 10^{-4}$	$1.3 \times 10^{-4}$	-
$ \lambda_{233} $	$5.6 \times 10^{-4}$	$3.1 \times 10^{-4}$	-
$ \lambda' $	$6.8 \times 10^{-5}$	$4.1 \times 10^{-5}$	-
$ \mu_e $ (GeV)	$8.1 \times 10^{-5}$	$3.4 \times 10^{-4}$	-
$ \mu_\mu $ (GeV)	$1.3 \times 10^{-4}$	$1.3 \times 10^{-4}$	-
$ \mu_\tau $ (GeV)	$1.3 \times 10^{-4}$	$1.3 \times 10^{-4}$	-

TABLE VIII. Bounds on the couplings for subcase 5 that satisfy MSW-LMA or MSW-SMA or vacuum, SuperK, CHOOZ and the neutrinoless beta decay constraints simultaneously.

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